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Effects of new physics in the rare decays $B \to K\ell^+\ell^-$ and $B \to K^*\ell^+\ell^-$ [1]

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Abstract

We parametrize phenomenologically possible new physics effects and calculate their influence on the invariant dilepton mass spectrum and the Dalitz plot for the decays $B \to K\ell^+\ell^-$ and $B \to K^*\ell^+\ell^-$. Especially the decay into K^* yields a wealth of new information on the form of the new interactions since the Dalitz plot is sensitive to subtle interference effects. We also show how transversely polarized K^*-mesons give additional information.

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1 Introduction

Rare decays of B-mesons, such as the recently observed processes $B \to K^*\gamma$ and $B \to X_s\gamma$ [1, 2] may become an important tool for studying new forces beyond the standard model. Indeed, many authors have investigated the effects of multi-Higgs models, supersymmetric theories, left-right symmetric models, etc. on this process [3]. The interest in these decays stems from the fact that they occur in the standard model only through loops and are therefore particularly sensitive to "new physics".

Whereas most of the past work has been concernend with the decay $b \to s\gamma$, the leptonic transition $b \to s\ell^+\ell^-$ is more sensitive to the actual form of the new interactions since it not only allows to measure a total rate, but also various kinematical distributions. More specifically, while new forces will only change the normalization of $b \to s\gamma$ (or $B \to K^*\gamma$), the interplay of the various operators will also change the spectra in the decays with two leptons. Thus, different new physics models which yield the same $b \to s\gamma$ amplitude may be distinguished by their characteristic invariant mass spectrum of the lepton pair in $b \to s\ell^+\ell^-$. Of course, there is a spectrum for the photon in the inclusive process $B \to X_s\gamma$, but it is essentially due to the wave function of the decaying meson and therefore not of interest here.

In this note we will parametrize possible new physics in a phenomenological way and evaluate their effects on the exclusive decays $B \to K\ell^+\ell^-$ and $B \to K^*\ell^+\ell^-$. Similar results also follow for the inclusive lepton pair spectrum, but since coherence effects tend to be more pronounced in exclusive decays we have concentrated here on these. Of course, this introduces a certain dependence on the hadronic model used. However, since most models do not disagree violently in the regions of interest and the crucial signs between the form factors are quite safe, we feel that a reliable representation of the effects is indeed possible. We have made use of a recent quark model [4, 5, 6] which gives rather satisfactory results. Needless to say that it reproduces the heavy quark [7] limit.

2 Effective Hamiltonian

The effects of both the new and the heavy standard model particles can be summarized by an effective Hamiltonian; for $b \to s\ell^+\ell^-$ (and the corresponding exclusive decays) it has the form [8]

$$\mathcal{H} = \frac{4G_F\lambda_t}{\sqrt{2}} \sum_{i=1}^{10} C_i(\mu)\, O_i(\mu) \quad , \tag{1}$$

where G_F denotes the Fermi constant and $\lambda_t = V_{tb}V_{ts}^*$ are products of CKM matrix elements. The $C_i(\mu)$ are coefficients calculated from processes with heavy particles exchanged and O_i are the following operators

$$\begin{aligned} O_1 &= (\bar{s}_\alpha\gamma^\mu L b_\alpha)\ (\bar{c}_\beta\gamma_\mu L c_\beta) \quad , \\ O_2 &= (\bar{s}_\alpha\gamma^\mu L b_\beta)\ (\bar{c}_\beta\gamma_\mu L c_\alpha) \quad , \end{aligned}$$

$$
\begin{aligned}
O_7 &= \frac{e}{16\pi^2}\, m_b\, (\bar{s}_\alpha \sigma^{\mu\nu} R b_\alpha)\, F_{\mu\nu} \quad , \\
O_{7'} &= \frac{e}{16\pi^2}\, m_b\, (\bar{s}_\alpha \sigma^{\mu\nu} L b_\alpha)\, F_{\mu\nu} \quad , \\
O_8 &= \frac{e^2}{16\pi^2}\, (\bar{s}_\alpha \gamma^\mu L b_\alpha) \bar{e}\gamma_\mu e \quad , \\
O_{8'} &= \frac{e^2}{16\pi^2}\, (\bar{s}_\alpha \gamma^\mu R b_\alpha) \bar{e}\gamma_\mu e \quad , \\
O_9 &= \frac{e^2}{16\pi^2}\, (\bar{s}_\alpha \gamma^\mu L b_\alpha) \bar{e}\gamma_\mu \gamma_5 e \quad , \\
O_{9'} &= \frac{e^2}{16\pi^2}\, (\bar{s}_\alpha \gamma^\mu R b_\alpha) \bar{e}\gamma_\mu \gamma_5 e \quad .
\end{aligned}
\tag{2}
$$

The basis of ref. [8] has been supplemented by the primed operators; the operators not written give small contributions. In eq. (1), the renormalization scale μ is usually chosen to be $\mu \simeq m_b$ in order to avoid large logarithms in the matrix elements of the operators O_i. Accordingly, the C_i are shifted from their perturbative value at $\mu \simeq m_W$ by the renormalization group equations. As the analytic expressions for the coefficients $C_i(\mu)$ in the standard model are given in many places in the literature [9], we only note here their numerical values which correspond to $m_t = 174$ GeV [10]: At the scale $\mu = m_W$ one gets

$$
\begin{aligned}
C_1(m_W) &= 0 \ , \ C_2(m_W) = -1 \ , \ C_7(m_W) = 0.196 \\
C_{7'}(m_W) &= \frac{m_s}{m_b} C_7(m_W) \ , \ C_8(m_W) = -2.02 \\
C_9(m_W) &= 4.56 \ , \ C_{8'}(m_W) = C_{9'}(m_W) = 0 \quad ,
\end{aligned}
\tag{3}
$$

while at $\mu = m_b$ the non-vanishing coefficients are

$$
\begin{aligned}
C_1(m_b) &= 0.226 \ , \ C_2(m_b) = -1.096 \ , \ C_7(m_b) = 0.310 \\
C_{7'}(m_b) &= \frac{m_s}{m_b} C_7(m_b) \ , \ C_8(m_b) = -3.84 \ , \ C_9(m_b) = 4.56 \quad .
\end{aligned}
\tag{4}
$$

It is well-known that the operators O_1 and O_2 generate both short and long distance contributions to $b \to s\ell^+\ell^-$ and to the corresponding exclusive decays. The long distance contributions are mainly due to the J/ψ and ψ' resonances [11, 12, 13]. These short and long distance contributions can be easily built-in [8, 14] by replacing the coefficient C_8 by C_8^{eff}, i.e.

$$
\begin{aligned}
C_8^{eff}(m_b) &= C_8(m_b) + [3C_1(m_b) + C_2(m_b)] \times \\
&\times \left[h(\hat{m}_c, \hat{s}) - \frac{3}{\alpha_{em}^2} \kappa \sum_{V_i = J/\psi, \psi'} \frac{\pi \Gamma(V_i \to \ell\ell) M_{V_i}}{q^2 - M_{V_i} + i M_{V_i} \Gamma_{V_i}} \right] \quad ,
\end{aligned}
\tag{5}
$$

where $\hat{m}_c = m_c/m_b$ and $\hat{s} = q^2/m_b^2$ with q^2 beeing the invariant mass squared of the lepton pair. The short distance contributions are contained in the function $h(\hat{m}_c, \hat{s})$

which reflects the one-loop matrix element of the four-quark operators O_1 and O_2. For $\hat{s} > 4z^2$

$$
\begin{aligned}
h(z,\hat{s}) &= -\left\{\frac{4}{9}\log z^2 - \frac{8}{27} - \frac{16}{9}\frac{z^2}{\hat{s}} + \frac{2}{9}\sqrt{1-\frac{4z^2}{\hat{s}}}\left(2+\frac{4z^2}{\hat{s}}\right)\right. \\
&\quad \times \left.\left[\log\left|\frac{1+\sqrt{1-4z^2/\hat{s}}}{1-\sqrt{1-4z^2/\hat{s}}}\right| + i\pi\right]\right\}
\end{aligned}
\tag{6}
$$

and for $\hat{s} < 4z^2$

$$
\begin{aligned}
h(z,\hat{s}) &= -\left\{\frac{4}{9}\log z^2 - \frac{8}{27} - \frac{16}{9}\frac{z^2}{\hat{s}} + \right. \\
&\quad \left. +\frac{4}{9}\sqrt{\frac{4z^2}{\hat{s}}-1}\left(2+\frac{4z^2}{\hat{s}}\right)\operatorname{atan}\left(\frac{1}{\sqrt{4z^2/\hat{s}-1}}\right)\right\} \quad .
\end{aligned}
\tag{7}
$$

The value of κ in eq. (5) must be chosen such that the combination

$$\kappa\,[3C_1(m_b) + C_2(m_b)] \approx -1$$

in order to correctly reproduce the branching ratio [14]

$$\mathrm{BR}(B \to J/\psi X \to \ell^+\ell^- X) = \mathrm{BR}(B \to J/\psi X)\ \mathrm{BR}(J/\psi \to \ell^+\ell^-) \quad .$$

In this language, the physics beyond the standard model is characterized by the values of the Wilson coefficients C_i at the perturbative scale m_W. For instance, supersymmetry can lead to a C_7 which is opposite to its standard model value [15], whereas left-right symmetric models could give a range for $C_{7'}$ between minus and plus the standard model value of C_7 [3, 16, 17]. Models with leptoquarks may cause the other coefficients to lie within similar intervals. In order to limit ourselves to a reasonable set of values in the examples, we chose to vary the coefficients between minus and plus their standard model values. In addition, we require that the rates for $B \to X_s\gamma$ and $B \to K^*\gamma$ do not change. For both reactions, this implies that we must leave the value of $(|C_7(m_b)|^2 + |C_{7'}(m_b)|^2)$ fixed within certain limits.

The running of the coefficients $C_{7'}$, $C_{8'}$ and $C_{9'}$ is simpler than the one of the unprimed coefficients since there is no operator mixing with O_2. In principle, new operators $O_{1'}$ and $O_{2'}$ with right-handed four fermion interactions should be included; we assume them to be small, as indicated by the high accuracy of the $V-A$ form in many weak decays, and we neglect them. For the scaling we can simply write in leading-log accuracy

$$
\begin{aligned}
C_{8'}(m_b) &= C_{8'}(m_W) \\
C_{9'}(m_b) &= C_{9'}(m_W) \\
C_{7'}(m_b) &= \eta^{-16/23} C_{7'}(m_W) \quad ,
\end{aligned}
\tag{8}
$$

where $\eta = \alpha_s(m_b)/\alpha_s(m_W) \simeq 1.71$ for $m_b = 5$ GeV.

3 The decay $B \to K\ell^+\ell^-$

The relevant hadronic matrix elements for the decay $B \to K\ell^+\ell^-$ associated with the operators O_i are parametrized as in eqs. (2.11) and (2.12) of ref. [5], i.e.

$$\langle p_K|\bar{s}\gamma_\mu(1 \mp \gamma_5)b|p_B\rangle = F_+(q^2)P_\mu + F_-(q^2)q_\mu \quad , \tag{9}$$

$$\langle p_K|\bar{s}i\sigma_{\mu\nu}q^\nu(1 \mp \gamma_5)b|p_B\rangle = \frac{1}{m_B + m_K}\left[P_\mu q^2 - (m_B^2 - m_K^2)q_\mu\right] F_T(q^2) \quad , \tag{10}$$

where $P = p_B + p_K$ and $q = p_B - p_K$ and $\sigma_{\mu\nu}$ is defined as $\sigma_{\mu\nu} = (i/2)[\gamma_\mu, \gamma_\nu]$. The form factors F_+, F_- and F_T are written generically as

$$F(q^2) = \frac{F(0)}{1 - q^2/\Lambda_1^2 + q^4/\Lambda_2^4} \quad . \tag{11}$$

The parameters Λ_1 and Λ_2 characterizing the shape of each of the form factors are given together with $F(0)$ in table 1 which has been taken from ref. [5]. The full information on the decay depends on two kinematical variables which we choose to be $\hat{s} = q^2/m_B^2$ and $x = E_\ell/m_B$, where q^2 is the invariant mass of the lepton pair and E_ℓ is the energy of the negatively charged lepton measured in the rest frame of the B meson. Using the notation of ref. [19] we get the double differential decay width

$$\frac{d^2\Gamma(B \to K\ell^+\ell^-)}{dx d\hat{s}} = \frac{G_F^2 m_B^5 |\lambda_t|^2}{8\pi^3} \frac{\alpha_{em}^2}{16\pi^2}\, \beta \left[4x(x_m - x) - \hat{s}(1 - 2x)\right] \quad , \tag{12}$$

where

$$\beta = \left|(C_8^{eff} + C_{8'})F_+ - \frac{2m_b(C_7 + C_{7'})}{m_B + m_K} F_T\right|^2 + |(C_9 + C_{9'})F_+|^2 \quad . \tag{13}$$

The variables x and $\hat{s}$ vary in the range

$$\frac{4m_\ell^2}{m_B^2} \le \hat{s} \le \frac{(m_B - m_K)^2}{m_B^2} \quad , \quad \frac{\hat{s} + 2x_m - \sqrt{\phi}}{4} \le x \le \frac{\hat{s} + 2x_m + \sqrt{\phi}}{4} \quad , \tag{14}$$

with

$$x_m = \frac{m_B^2 - m_K^2}{2m_B^2} \quad , \quad \phi = (\hat{s} + 2x_m)^2 - 4\hat{s}. \tag{15}$$

Integrating eq. (12) over the variable x in the range specified in eq. (14) yields the invariant mass distribution of the lepton pair

$$\frac{d\Gamma(B \to K\ell^+\ell^-)}{d\hat{s}} = \frac{G_F^2 m_B^5 |\lambda_t|^2}{96\pi^3} \frac{\alpha_{em}^2}{16\pi^2}\, \beta\, \phi^{3/2} \quad . \tag{16}$$

We notice the relative sign in eq. (13) between the terms involving the C_7 and C_8 coefficients which leads to a potentially interesting interference effect. Since F_+ and F_T are calculated in a specific model, one might worry that the sign is strongly

model dependent. However, in the approximation where the quarks are on-shell, the Gordon decomposition confirms the sign.

In Fig. 1 we show the dilepton invariant mass distribution for the standard model and for a model where C_8 has the opposite sign from the standard model value. The difference between the two curves is especially noticeable near the resonances, a fact which is easily explained by eq. (5) (see also eq. (13)). On the other hand, changing the sign of C_7 does not alter the standard model picture very much, since the contributions from the C_7 terms in eq. (13) are small. From the form of β we see that other variations of the coefficients will not give dramatic changes of the spectra and Dalitz plot.

In the numerical evaluations we take $m_b = m_B = 5.28$ GeV in all formulas except in the Wilson coefficients, where we used $m_b = 5$ GeV. Furthermore, in all the plots we divided the spectra by the total B-meson decay width, estimated to be [18]

$$\Gamma_{tot} = \frac{f\, m_b^5 G_F^2 |V_{cb}|^2}{192\pi^3} \quad , \tag{17}$$

with $f \approx 3$.

4 The decay $B \to K^* \ell^+ \ell^-$

Next, we consider the more interesting decay $B \to K^* \ell^+ \ell^-$. Again, the relevant hadronic matrix elements are parametrized as in eqs. (2.18) and (2.19) of ref. [5], i.e.

$$\begin{aligned}\langle p_{K^*} | \bar{s} \gamma_\mu (1 \mp \gamma_5) b | p_B \rangle = \frac{1}{m_B + m_{K^*}} & \Big[-iV(q^2) \epsilon_{\mu\nu\alpha\beta} \epsilon^{*\nu} P^\alpha q^\beta \\ \pm A_0(q^2)(m_B^2 - m_{K^*}^2)\epsilon^*_\mu \pm A_+(q^2)(\epsilon^* P) P_\mu & \pm A_-(q^2)(\epsilon^* P) q_\mu \Big]\end{aligned} \tag{18}$$

$$\begin{aligned}&\langle p_{K^*} | \bar{s} i \sigma_{\mu\nu} q^\nu (1 \pm \gamma_5) b | p_B \rangle = -ig(q^2) \epsilon_{\mu\nu\alpha\beta} \epsilon^{*\nu} P^\alpha q^\beta \pm \\ &a_0(q^2)\,(m_B^2 - m_K^2) \left[\epsilon^*_\mu - \frac{1}{q^2} (\epsilon^* q) q_\mu \right] \pm a_+(q^2)\,(\epsilon^* P) \left[P_\mu - \frac{1}{q^2}(Pq) q_\mu \right] \quad ,\end{aligned} \tag{19}$$

where P and q are now defined as $P = p_B + p_{K^*}$ and $q = p_B - p_{K^*}$. The ϵ-tensor is taken in the Bjorken-Drell convention (i.e. $\epsilon_{0123} = 1$) contrary to ref. [5], which explains the additional minus sign of the terms proportional to the ϵ-tensor. The various form factors in eqs. (18) and (19) are written as in eq. (11) and the parameters are given in table 1. With the notation of ref. [19] we get

$$\begin{aligned}\frac{d^2\Gamma(B \to K^* \ell^+ \ell^-)}{dx d\hat{s}} \quad = \quad & \frac{G_F^2 m_B^5 |\lambda_t|^2}{4\pi^3} \frac{\alpha_{em}^2}{16\pi^2} \left[\frac{\hat{s}}{m_B^2}\, \alpha + 2\beta[4x(x_m - x) - \hat{s}(1-2x)] \right. \\ & \left. -2\gamma \hat{s}(x_m - 2x + \hat{s}/2) \right] \quad ,\end{aligned} \tag{20}$$

where

$$
\begin{aligned}
\alpha &= |G_A^0|^2 + |F_A^0|^2 + 4m_B^2 p_s^2 \left(|G_V|^2 + |F_V|^2\right) \\
\beta &= \frac{|G_A^0|^2 + |F_A^0|^2}{4m_{K^*}^2} - m_B^2 \hat{s} \left(|G_V|^2 + |F_V|^2\right) + \frac{m_B^2}{m_{K^*}^2} p_s^2 \left(|G_A^+|^2 + |F_A^+|^2\right) + \\
&\quad \frac{1}{2}\left[\frac{m_B^2}{m_{K^*}^2}(1-\hat{s}) - 1\right] \mathrm{Re}(G_A^0 G_A^+ + F_A^0 F_A^+) \\
\gamma &= -2\mathrm{Re}(G_A^0 F_V + F_A^0 G_V) \quad ,
\end{aligned}
\tag{21}
$$

and

$$
p_s^2 = \phi \frac{m_B^2}{4} \quad .
$$

The functions G_V, F_V, G_A^0, F_A^0, G_A^+ and F_A^+ are combinations of form factors and Wilson coefficents; they read

$$
\begin{aligned}
G_V &= \frac{(C_8^{eff} + C_{8'})V}{2(m_B + m_{K^*})} - (C_7 + C_{7'}) \frac{m_b}{q^2} g \\
F_V &= \frac{(C_9 + C_{9'})V}{2(m_B + m_{K^*})} \\
G_A^0 &= \frac{(C_8^{eff} - C_{8'})(m_B - m_{K^*})A_0}{2} - (C_7 - C_{7'}) \frac{m_b}{q^2} (m_B^2 - m_{K^*}^2) a_0 \\
F_A^0 &= \frac{(C_9 - C_{9'})}{2} (m_B - m_{K^*}) A_0 \\
G_A^+ &= \frac{(C_8^{eff} - C_{8'})A_+}{2(m_B + m_{K^*})} - (C_7 - C_{7'}) \frac{m_b}{q^2} a_+ \\
F_A^+ &= \frac{(C_9 - C_{9'})A_+}{2(m_B + m_{K^*})} \quad .
\end{aligned}
\tag{22}
$$

The range of the variables x and $\hat{s}$ and the definition of x_m are obtained from eqs. (14) and (15) by the obvious replacement $m_K \to m_{K^*}$.

Integration over the variable x leads to the invariant mass distribution of the lepton pair

$$
\frac{d\Gamma(B \to K^* \ell^+ \ell^-)}{d\hat{s}} = \frac{G_F^2 m_B^5 |\lambda_t|^2}{8\pi^3} \frac{\alpha_{em}^2}{16\pi^2} \sqrt{\phi} \left[\frac{\hat{s}}{m_B^2} \alpha + \frac{\beta}{3} \phi\right] \quad . \tag{23}
$$

Taking into account only transversely polarized K^* mesons in the final state, the analogous expression for the double differential decay width is

$$
\begin{aligned}
\frac{d^2\Gamma^T(B \to K^* \ell^+ \ell^-)}{dx d\hat{s}} &= \frac{G_F^2 m_B^5 |\lambda_t|^2}{4\pi^3} \frac{\alpha_{em}^2}{16\pi^2} \left[\frac{\hat{s}}{m_B^2} \alpha^T + 2\beta^T [4x(x_m - x) - \hat{s}(1-2x)] \right. \\
&\left. -2\gamma^T \hat{s}(x_m - 2x + \hat{s}/2) - \frac{\delta^T}{2m_B^2} \hat{s} \left(\frac{4x(x_m - x) - \hat{s}(1-2x)}{(x_m + \hat{s}/2)^2 - \hat{s}}\right)\right] \quad ,
\end{aligned}
\tag{24}
$$

where

$$\alpha^T = \alpha \ , \ \beta^T = -m_B^2 \hat{s}(|G_V|^2 + |F_V|^2) \ , \ \gamma^T = \gamma \ , \ \delta^T = |G_A^0|^2 + |F_A^0|^2 \ . \tag{25}$$

The corresponding invariant mass distribution of the lepton pair reads

$$\begin{aligned} \frac{d\Gamma^T(B \to K^* \ell^+ \ell^-)}{d\hat{s}} &= \frac{G_F^2 m_B^5 |\lambda_t|^2}{8\pi^3} \frac{\alpha_{em}^2}{16\pi^2} \sqrt{\phi} \left[\frac{\hat{s}}{m_B^2} \alpha^T + \frac{\beta^T}{3} \phi - \right. \\ & \left. \frac{\delta^T}{12 m_B^2} \hat{s} \frac{\phi}{(x_m + \hat{s}/2)^2 - \hat{s}} \right] \ . \end{aligned} \tag{26}$$

In Figs. 2 to 4 we show the invariant mass distribution and the Dalitz plots for three sets of values of the couplings, all of which give the same rate for the decay $B \to X_s \gamma$. In Fig. 2 we illustrate the standard model prediction (Fig. 2a) and the case where the sign of C_8 is reversed (Fig. 2b); this again leads to the marked deviation around the resonances. In addition, we see from eqs. (21) and (23) that reversing the sign of C_9 does not change the invariant mass distribution; however, this change influences the Dalitz plot significantly as illustrated in Fig 3: While the standard model distribution in Fig. 3a is populated more for large x-values, Fig. 3c shows more points for small values of x.

Finally, from Fig. 4 we see that the transverely polarized K^* mesons accentuate this asymmetry in the Dalitz plot and can therefore be used to fix the sign of the coefficient C_9.

5 Conclusions

In this paper we have investigated the effects of new physics on the rare decays $B \to K \ell^+ \ell^-$ and $B \to K^* \ell^+ \ell^-$. Whereas the new physics influence only the rate in the decay $B \to X_s \gamma$, they change significantly the various spectra in the leptonic decays. Thus, while certain models give the same rate for $B \to X_s \gamma$, they can be distinguished by their dilepton spectra. In addition, the Dalitz plot carries further information (compare Figs. 3 and 4). Furthermore, looking only at the transversely polarized K^* yields improved information; as an example we have investigated the coefficient C_9.

Of course, the applicability of our results will depend on the number of avaliable decays. It is expected that at the hadronic B facilities at LHC the required number (about 100 for the dilepton spectrum and 1000 for the Dalitz plot) can be analyzed. Then, these decays will be a truly powerful tool towards understanding new physics.

Noted added: A. Ali, G. Giudice and T. Mannel have investigated in a similar spirit the inclusive dilepton decays (private communication and [20]).

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Figure captions

Fig. 1: Invariant dilepton mass distribution for the decay $B \to K\ell^+\ell^-$. In Fig. 1a the prediction within the standard model is given; Fig. 1b shows the result for a model where $C_8(m_W) = 2.02$ (opposite value of the standard model; see eq. (3)) and the other coefficients unchanged.

Fig. 2: Same as Fig. 1, but for the decay $B \to K^*\ell^+\ell^-$. The solid (dashed) line takes into account all (only the transverse) polarizations of the K^*-meson.

Fig. 3: Dalitz plot (compare eq. (20)) for the decay $B \to K^*\ell^+\ell^-$ taking into account all polarizations of the K^*-meson. We have used 8000 points in the scatter plots. The distribution in Fig. 3a is for the standard model values of the coefficients while Fig. 3b has the sign of $C_8(m_W)$ reversed. In Fig. 3c we used the standard model values of the the coefficients, except that the sign of $C_9(m_W)$ has been reversed.

Fig. 4: Dalitz plot for transversely polarized K^*-mesons. Fig. 4a is for the standard model values of the coefficients and Fig. 4b has the sign of $C_9(m_W)$ reversed.

Tables

	F_+	F_T	V	A_+	A_0	g	a_+	a_0
$F(0)$	0.30	-0.30	-0.35	0.24	-0.37	0.31	-0.31	0.31
Λ_1 (GeV)	4.06	4.08	4.05	4.39	5.92	4.05	4.37	6.40
Λ_2 (GeV)	5.51	5.50	5.45	5.81	8.54	5.44	5.84	8.84

Table 1: Values of the various $F(0)$, Λ_1 and Λ_2